

# On the derivation of an effective Higgs field

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In one respect, the massive vector-boson shows its difference from a massless vector-boson by one more physical polarization, known as longitudinal polarization. In another respect, the quantized boson acquires its mass by Higgs mechanism. In this paper we study the effect of the longitudinal polarization in  $U(1)$  case by substituting it into the primary Yang-Mills Lagrangian  $-\frac{1}{4}F_{\mu\nu}F^{\mu\nu}$ . Under a hypothesis of strong transversal condition for free vector boson, it is found that in the Lagrangian the scalar field for the Higgs mechanism can automatically arise after we separate a part equivalent to the contribution of a massless boson. In addition, a criterion is obtained to infer whether the boson is massive or not: if  $\mathbf{E}^2 - \mathbf{B}^2 \neq 0$ , where  $\mathbf{E}$  and  $\mathbf{B}$  are field strengths, then it is massive. The analysis also pertains to  $SU(2)$  case. The method in this paper is performed before any quantizations.

The newly designed tera-scale Large Hadron Collider (LHC) of CERN is expected to begin its switch within weeks [1] [2]. One of its missions is to search for Higgs particle. The "unexpected" results, which push the possibility of the existence of Higgs particle to higher energy scale, may appear at LHC. Then the conventional Higgs mechanism [3]—which was thought to be accompanied by one light Higgs particle—has to be modified. However, the light Higgs particle ( $m_h \lesssim 144\text{GeV}$ ) meets the stringent constraints by electroweak gauge couplings [4], and the most recent experiments involving these constraints also point to a low value of Higgs mass [5].

On the theoretical side, the Higgs mechanism is made a cornerstone related to the vacuum in constructing a renormalizable theory with massive vector bosons. It provides delicate cancellations of  $\xi$ -dependent (gauge-fixing-condition dependent) terms in calculating  $S$ -matrix elements of all orders. In spite of the success of Higgs mechanism, Veltman [6] insists that the scalar field required by a renormalizable theory should have other origins if the light Higgs particle were not found. To seek another scheme in place of Higgs mechanism, however, is not feasible at present stage. Here we find a way out to validate Higgs mechanism without the need of material Higgs field, in case that the light Higgs particle is missing in experiment. The goal of this paper is to separate a term  $\partial_\mu\varphi\partial^\mu\varphi$  for vacuum states from the free-field Lagrangian density  $-\frac{1}{4}F_{\mu\nu}F^{\mu\nu}$  by substituting the physical polarization vectors of massive boson.

The essence of Higgs mechanism includes two parts: one is the hypothesis that there exists a scalar field responsible for the vacuum state (this paper we name it **vacuum field**), which couples with vector bosons via covariant derivative,  $\partial_\mu \rightarrow \partial_\mu + ieA_\mu$ ; the other is, e.g. in  $U(1)$  case, the specification of self-interaction of the scalar field, which makes the global symmetry spontaneously broken and meanwhile the scalar field gets a vacuum expectation value  $\phi_0$ . Subtracted by the expectation value, the real part of the scalar field is defined as the Higgs field. The complete form of Lagrangian involving Higgs mechanism reads

$$\mathcal{L} = -\frac{1}{4}(F_{\mu\nu})_{\text{massless}}^2 + |D_\mu\phi|^2 - V(\phi), \quad (1)$$

in which  $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$  and  $D_\mu = \partial_\mu + ieA_\mu$ . For simplicity the subscript "massless" means only two physical (spatially transversal) polarizations of massless bosons are present in  $A_\mu$ . The conventional decomposition of  $\phi(x)$  is always used,

$$\phi(x) = \phi_0 + \frac{1}{\sqrt{2}}(\phi_1(x) + i\phi_2(x)), \quad (2)$$

where  $\phi_1$  is Higgs field and  $\phi_2$  is Goldstone field. In what follows we mainly focus on the derivation of the term  $|D_\mu\phi|^2$  in Eq. (1) by substituting the massive polarization vectors into the primary Yang-Mills Lagrangian  $-\frac{1}{4}F_{\mu\nu}F^{\mu\nu}$ .

In  $U(1)$  case and for a massive neutral spin-1 vector boson which freely moves along the third direction (later labelled as  $z$  direction), there exists a particular frame of reference where the three polarizations have the form [7]

$$\begin{aligned} \epsilon_1(\mathbf{k}) &= (0, 1, 0, 0) \\ \epsilon_2(\mathbf{k}) &= (0, 0, 1, 0) \\ \epsilon_3(\mathbf{k}) &= (|\mathbf{k}|, 0, 0, E_{\mathbf{k}})/m, \end{aligned} \quad (3)$$

here  $k = (E_{\mathbf{k}}, 0, 0, |\mathbf{k}|)$  and  $k^2 = k_\mu k^\mu = m^2$ . The polarization states meet the requirements of  $k_\mu \epsilon(\lambda)^\mu(\mathbf{k}) = 0$  and  $\epsilon_{\lambda\mu}(\mathbf{k}) \epsilon_{\lambda'\mu}^\mu(\mathbf{k}) = g_{\lambda\lambda'}$ , here the summation convention (repeated Greek indices are summed) is understood and the nonzero components of metric tensor  $g_{\lambda\lambda'}$  are  $g_{00} = -g_{11} = -g_{22} = -g_{33} = +1$ . Here the temporal polarization is absent since for massive bosons only these three are "physical" [7].

The polarization vectors  $\epsilon_\lambda(\mathbf{k})$  and the four-vector potential  $A(t, \mathbf{x})$  ( $\mathbf{x} = (x, y, z)$ ) are linked by the Fourier expansion. We simplify the expansion by the above chosen frame of reference, which is arbitrary but remains fixed temporarily. In this frame, the plane wave has momentum  $\mathbf{k}$ ,

$$\begin{aligned} A(t, \mathbf{x}) &= \int \frac{d\mathbf{k}}{\sqrt{2\omega_{\mathbf{k}}(2\pi)^3}} \sum_{\lambda=1}^3 \epsilon_\lambda(\mathbf{k}) (a_\lambda(\mathbf{k}) e^{-ik \cdot x} + a_\lambda^*(\mathbf{k}) e^{ik \cdot x}) \\ &= A_1(t, \mathbf{x}, \mathbf{k}) \epsilon_1(\mathbf{k}) + A_2(t, \mathbf{x}, \mathbf{k}) \epsilon_2(\mathbf{k}) + A_3(t, \mathbf{x}, \mathbf{k}) \epsilon_3(\mathbf{k}), \end{aligned} \quad (4)$$

where  $\omega_{\mathbf{k}} = (m^2 + \mathbf{k}^2)^{1/2}$ . We will substitute the Eq. (4) into the Lagrangian  $-\frac{1}{4}F_{\mu\nu}F^{\mu\nu}$  to compare the result with the case when two transversal polarization vectors  $\epsilon_1(\mathbf{k})$  and  $\epsilon_2(\mathbf{k})$  [These two are physical polarizations for photon] are present only, i.e. to find the effect of  $\epsilon_3(\mathbf{k})$  in  $-\frac{1}{4}F_{\mu\nu}F^{\mu\nu}$ .

Since the Lagrangian  $-\frac{1}{4}F_{\mu\nu}F^{\mu\nu}$  is a Lorentz invariant quantity, it doesn't matter if we perform a Lorentz transformation on four-vector potential  $A_\mu(t, \mathbf{x})$ . In view of the last line of Eq. (4), the transformation will appear only to  $\epsilon_1(\mathbf{k})$ ,  $\epsilon_2(\mathbf{k})$  and  $\epsilon_3(\mathbf{k})$ , and meanwhile the values of coefficients  $A_\lambda(t, \mathbf{x}, \mathbf{k})$  will vary with  $\mathbf{k}$ , explicitly as  $\frac{(a_\lambda(\mathbf{k})e^{-ik \cdot x} + a_\lambda^*(\mathbf{k})e^{ik \cdot x})}{\sqrt{2\omega_{\mathbf{k}}}}$ . Let's suppose the system is boosted along the  $z$  direction with the velocity  $v_z = \frac{|\mathbf{k}|}{E_{\mathbf{k}}}$  (in natural unit, light velocity  $c = 1$ ) [8]. This is possible because the vector  $\epsilon_3(\mathbf{k})$  is space-like ( $\frac{E_{\mathbf{k}}}{m})^2 - (\frac{|\mathbf{k}|}{m})^2 = 1 > 0$ ). Then in the new reference frame

$$|\mathbf{k}'| = \frac{|\mathbf{k}| - v_z E_{\mathbf{k}}}{\sqrt{1 - v_z^2}}, \quad E'_{\mathbf{k}} = \frac{E_{\mathbf{k}} - v_z |\mathbf{k}|}{\sqrt{1 - v_z^2}}. \quad (5)$$

Note that  $|\mathbf{k}| = v_z E_{\mathbf{k}}$ , one obtains

$$E'_{\mathbf{k}} = \frac{E_{\mathbf{k}} - v_z^2 E_{\mathbf{k}}}{\sqrt{1 - v_z^2}} = \sqrt{E_{\mathbf{k}}^2 - |\mathbf{k}|^2} = m, \quad (6)$$

and thus

$$\epsilon'_L(\mathbf{k}) = (0, 0, 0, 1). \quad (7)$$

Whereas under this transformation the other two polarization vectors  $\epsilon_1(\mathbf{k})$  and  $\epsilon_2(\mathbf{k})$  remain unchanged. In view of Eq. (7) and the fact that massless vector boson owns only two physical polarizations, one concludes that **the massless vector fields are not the limit of zero mass of massive vector fields** (which is also mentioned in another manner in [6]), since in Eq. (7) the longitudinal polarization is mass independent. The potential  $A_\mu$  in Eq. (4) now yields

$$A'(t, \mathbf{x}) = A'_1(t, \mathbf{x}, \mathbf{k}) \epsilon_1(\mathbf{k}) + A'_2(t, \mathbf{x}, \mathbf{k}) \epsilon_2(\mathbf{k}) + A'_3(t, \mathbf{x}, \mathbf{k}) \epsilon'_L(\mathbf{k}). \quad (8)$$

Next let's examine how the third polarization (i.e.  $A_3$  or  $A_z$ ) in Eq.(7) affects the Lagrangian  $-\frac{1}{4}(F_{\mu\nu})^2$ . The  $U(1)$  field-strength tensor has the following form

$$\begin{aligned} F^{\mu\nu} &= \begin{pmatrix} 0 & \frac{\partial}{\partial x} A_t - \frac{\partial}{\partial t} A_x & \frac{\partial}{\partial y} A_t - \frac{\partial}{\partial t} A_y & \frac{\partial}{\partial z} A_t - \frac{\partial}{\partial t} A_z \\ -\frac{\partial}{\partial x} A_t + \frac{\partial}{\partial t} A_x & 0 & \frac{\partial}{\partial y} A_x - \frac{\partial}{\partial x} A_y & \frac{\partial}{\partial z} A_x - \frac{\partial}{\partial x} A_z \\ -\frac{\partial}{\partial y} A_t + \frac{\partial}{\partial t} A_y & -\frac{\partial}{\partial y} A_x + \frac{\partial}{\partial x} A_y & 0 & \frac{\partial}{\partial z} A_y - \frac{\partial}{\partial y} A_z \\ -\frac{\partial}{\partial z} A_t + \frac{\partial}{\partial t} A_z & -\frac{\partial}{\partial z} A_x + \frac{\partial}{\partial x} A_z & -\frac{\partial}{\partial z} A_y + \frac{\partial}{\partial y} A_z & 0 \end{pmatrix} \\ &= \begin{pmatrix} 0 & E_x & E_y & E_z \\ -E_x & 0 & B_z & -B_y \\ -E_y & -B_z & 0 & B_x \\ -E_z & B_y & -B_x & 0 \end{pmatrix}, \end{aligned} \quad (9)$$

in which  $\mu$  represents the lines and  $\nu$  denotes the columns. Substituting the transformed potential  $A'_\mu$  in Eq.(8) [In what follows we will still refer to  $A_\mu$  as  $A'_\mu$ , with Eq.(7) as the longitudinal polarization] into Eq.(9), it is found that the terms including  $A_z$  appear only in the last line and the last column of the matrix. Considering the following strong transversal condition for free vector boson

$$\frac{\partial}{\partial z} A_x = \frac{\partial}{\partial z} A_y = \frac{\partial}{\partial z} A_z = 0, \quad (10)$$

(where we introduce the condition  $\frac{\partial}{\partial z}A_x = \frac{\partial}{\partial z}A_y = 0$  phenomenologically and  $\frac{\partial}{\partial z}A_z = 0$  is in coincidence with the Coulomb gauge fixing condition  $\mathbf{k} \cdot \mathbf{A} = 0$ ). In developing this paper, we have done our best not to be involved in the using of gauge fixing conditions, which are closely related with the quantization of fields) and substituting the component form of field strengths  $E_z = -\frac{\partial}{\partial t}A_z$ ,  $B_y = \frac{\partial}{\partial x}A_z$ ,  $B_x = -\frac{\partial}{\partial y}A_z$  into  $F^{\mu\nu}$  and subsequently into  $-\frac{1}{4}(F_{\mu\nu})^2_{massive}$ , we obtain the relevant terms including  $A_z$  as follows

$$\delta\mathcal{L} = -\frac{1}{4}(F_{\mu\nu})^2_{massive}(\text{relevant parts}) = \frac{1}{2}(E_z^2 - B_y^2 - B_x^2) . \quad (11)$$

If we introduce a scalar field  $\varphi = \frac{1}{\sqrt{2}}A_z(t, \mathbf{x}, \mathbf{k})$ , then the terms in  $\delta\mathcal{L}$  can be straightforwardly written as

$$\delta\mathcal{L} = (\partial_t\varphi)^2 - (\partial_x\varphi)^2 - (\partial_y\varphi)^2 = \partial_\mu\varphi\partial^\mu\varphi . \quad (12)$$

By this way one obtains the term  $\partial_\mu\varphi\partial^\mu\varphi$  in the Lagrangian for Higgs mechanism. The separation of this term requires that the component  $A_z(t, \mathbf{x}, \mathbf{k})$  should not be transformed to zero by gauge transformation, therefore the separation procedure here is by no means gauge-invariant.

So far the original Lagrangian  $-\frac{1}{4}(F_{\mu\nu})^2_{massive}$  falls into two parts:

$$\mathcal{L}_{massive} = -\frac{1}{4}(F_{\mu\nu})^2_{massive} = \mathcal{L}_{massless} + \delta\mathcal{L}, \quad (13)$$

where the first term of right hand side is the same as that for massless photon, in which only two polarization vectors  $\epsilon_1(\mathbf{k})$  and  $\epsilon_2(\mathbf{k})$  are present. In general one puts another two polarization vectors (longitudinal and scalar) additional to them to form complete polarization states of massless bosons (then make it mixed by Lorentz transformations). In scattering processes the physical effects of redundant polarizations are cancelled [7]. Let's add a self-interaction term  $V(\varphi)$  [somehow produced by vacuum] to Eq.(13), then the original Lagrangian for massive boson becomes

$$\mathcal{L} = -\frac{1}{4}(F_{\mu\nu})^2_{massless} + \partial_\mu\varphi\partial^\mu\varphi - V(\varphi) . \quad (14)$$

We can see that the Lagrangian already takes the form that for Higgs mechanism in Eq. (1), with  $\varphi$  identical with  $\phi$ , if only we make replacement  $\partial_\mu \rightarrow \partial_\mu + ieA_\mu$  to gain the coupling between  $\varphi(x)$  and  $A_\mu$  as usual. To use the Higgs mechanism, the next step is to make the **vacuum field**  $\varphi$  complex to include a part responsible for the Goldstone field, as shown in Eq. (2). Finally we should turn to the application of Higgs mechanism and the corresponding quantization etc..

Noting the physical fact that the value  $\mathcal{L}_{massless} = \mathbf{E}^2 - \mathbf{B}^2 = 0$ , we recognize that in Eq. (13)  $\mathcal{L}_{massive}$  must be nontrivial, i.e.  $\mathcal{L}_{massive} = \tilde{\mathbf{E}}^2 - \tilde{\mathbf{B}}^2 \neq 0$ , otherwise the existence of  $\epsilon_3(\mathbf{k})$  becomes nonsense [as usual, here  $E^i = F^{i0}$  and  $B^i = -\epsilon^{ijk}F^{jk}$  as default form for the massless cases, with  $\tilde{\mathbf{E}}$  and  $\tilde{\mathbf{B}}$  the same form but for massive cases]. Moreover, one should note that by reversing the above derivation, the field  $\varphi(x)$  turns out to be a component of  $A_\mu$ , gained by the aid of Lorentz transformation, and dependent on the breaking of gauge transformation.

The above separation steps are applicable to  $SU(2)$  case too. The definition of  $F_{\mu\nu}$  for non-Abelian gauge fields is

$$F_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a + g\epsilon^{abc}A_\mu^b A_\nu^c, \quad (15)$$

with  $g$  as coupling constant and indices  $a$  corresponding to the generators of gauge group. By performing the gauge transformation

$$A_\mu^a \rightarrow A_\mu^a + \frac{1}{g}(\partial_\mu\alpha^a) + f^{abc}A_\mu^b\alpha^c, \\ (f^{abc} \text{ are structure constants for gauge group}) \quad (16)$$

the last term in Eq. (18) can be transformed away by gauge, and thus the massive Lagrangian  $-\frac{1}{4}F_{a\mu\nu}F^{a\mu\nu}$  reads [9]

$$-\frac{1}{4}F_{a\mu\nu}F^{a\mu\nu} \rightarrow -\frac{1}{4}\tilde{F}_{a\mu\nu}\tilde{F}^{a\mu\nu} \\ \text{with } \tilde{F}_{\mu\nu}^a = \partial_\mu A_\nu^a - \partial_\nu A_\mu^a, \quad (17)$$

which is analogous to the  $U(1)$  case. Now we can perform the similar separation Eq. (7)  $\sim$  Eq. (13) in  $U(1)$  case on Eq. (17), the difference is that now the field strengths  $E^i = \tilde{F}^{i0}$  and  $B^i = -\epsilon^{ijk}\tilde{F}^{jk}$  are all matrices. For  $SU(2)$  case

three scalar fields  $\varphi^a$  [ $a = 1, 2, 3$ ] are acquired. But to fit the conventional application, we assume they are equal. Again we note that this procedure is not gauge invariant due to the derivation from Eq. (16) to Eq. (17).

In the  $SU(2)$  case we note that the field  $\varphi$  is also a component of the field  $A_\mu$ ,

$$\varphi = \frac{1}{\sqrt{2}} A_z(t, \mathbf{x}, \mathbf{k}) , \quad (18)$$

as in  $U(1)$  case. However, since the field  $A_\mu$  is a matrix,  $A_\mu^a \tau^a$  ( $\tau^a = \frac{\sigma^a}{2}$ , with  $\sigma^a$  the Pauli matrices), when we use the covariant differential  $D_\mu = \partial_\mu + ieA_\mu$  instead of  $\partial_\mu$  to produce the coupling between  $\varphi$  and  $A_\mu$ , at the first sight it seems problematic to make  $\varphi$  act as a  $SU(2)$  scalar  $\begin{pmatrix} \eta_1 + i\eta_2 \\ v + \sigma(x) + i\eta_3 \end{pmatrix}$  in the sense of the  $SU(2) \otimes U(1)$  standard model. Whereas we find the effect of  $\varphi$  is identical with that of  $SU(2)$  scalar  $\begin{pmatrix} 0 \\ v \end{pmatrix}$  if one reduces the summation  $A_z^a \tau^a$  (over the indices  $a$ ) to a particular term, e. g.  $A_z^3 \tau^3$ . Let's Recall the mass-term production in Higgs mechanism for  $SU(2)$  case, where we replace the differential  $\partial_\mu$  in  $\partial_\mu \varphi \partial^\mu \varphi$  with covariant one  $D_\mu \varphi = (\partial_\mu - igA_\mu^a \tau^a) \varphi$ , and make  $\varphi = \begin{pmatrix} 0 \\ v \end{pmatrix}$ , then the relevant part of the Lagrangian yields [10]

$$\begin{aligned} \Delta \mathcal{L} &= \frac{1}{2} \varphi (g A_\mu^a \tau^a) (g A^{\mu b} \tau^b) \varphi \\ &= \frac{1}{2} (0, v) (g A_\mu^a \tau^a) (g A^{\mu b} \tau^b) \begin{pmatrix} 0 \\ v \end{pmatrix} \\ &= \frac{1}{2} \frac{v^2 g^2}{4} A_\mu A^\mu , \end{aligned} \quad (19)$$

by which  $A_\mu$  gains a mass  $\frac{gv}{2}$ . Now the similar expression exists even if the field  $\varphi$  takes a matrix form  $\frac{1}{\sqrt{2}} A_z^3 \tau^3$ , as follows

$$\begin{aligned} \Delta \mathcal{L} &= \frac{1}{2} \varphi (g A_\mu^a \tau^a) (g A^{\mu b} \tau^b) \varphi \\ &= \frac{1}{2} \frac{1}{\sqrt{2}} \frac{1}{2} \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix} A_z^3 (g A_\mu^a \tau^a) (g A^{\mu b} \tau^b) \frac{1}{\sqrt{2}} \frac{1}{2} \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix} A_z^3 \\ &= \frac{1}{64} (A_z^3)^2 g^2 \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix} [(A_\mu^1)^2 + (A_\mu^2)^2 + (A_\mu^3)^2] \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix} \\ &= \frac{1}{64} (A_z^3)^2 g^2 A_\mu A^\mu , \end{aligned} \quad (20)$$

by this way the field  $A_\mu$  gains a mass proportional to  $(A_z^3)^2 g^2$ . So far making the matrix form in place of the original  $SU(2)$  scalar field hasn't induced any problem in producing boson's mass. However, this scheme seems not appropriate to the  $SU(3)$  case since  $(g A_\mu^a \lambda^a) (g A^{\mu b} \lambda^b)$  ( $\lambda^a$  is Gell-Mann matrices) cannot be written as  $(A_\mu^1)^2 + (A_\mu^2)^2 + \dots + (A_\mu^8)^2$  without cross terms.

In summary, in this note we find the **vacuum field**, and thus the Higgs field can be separated from the Lagrangian for massive boson field, whence we may feel easy even if at LHC the Higgs particle were not claimed after the near-future collection of data. From the separation procedure, one notes that part of the spontaneous breaking of symmetries in conventional Higgs mechanism has been transferred to the breaking of  $U(1)$  or  $SU(2)$  gauge symmetries before using the Higgs mechanism. So the separation we performed is by no means a gauge invariant procedure. However it is feasible—we have found where the energy of the vacuum states possibly hides in the massive Lagrangian  $-\frac{1}{4}(F_{\mu\nu})_{massive}^2$ . The result of the decomposition of massive-boson's Lagrangian provides us a criterion to judge whether a boson owns mass or not: If  $\tilde{\mathbf{E}}^2 - \tilde{\mathbf{B}}^2 \neq 0$ , then it does. We choose the field strengths  $\tilde{\mathbf{E}}$  and  $\tilde{\mathbf{B}}$  to express this criterion for they are observables even in classical cases. Furthermore, we conclude from Eq. (20) that the boson's mass is born from the self-interaction. Throughout this paper we have worked in physical polarizations, regardless of the completeness of them. In addition, the steps of separation happens before any quantizations are performed, so we don't get involved in any particular gauge fixing conditions, as well as the corresponding propagator or other Green functions. Consequently, the field strengths  $\tilde{\mathbf{E}}$  and  $\tilde{\mathbf{B}}$  are not operators yet.

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